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Quark-Model Predictions for Reactions with Hyperon Beams*

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Quark model predictions are discussed for Primakoff excitation of hyperon resonances, total hyperon-nucleon cross sections and diffractive excitation. The U spin selection rule forbidding electromagnetic excitation of negatively charged decuplet resonances is shown to hold even in the presence of large SU(3) symmetry breaking. A new model for diffractive excitation is presented which suggests the existence of new hyperon resonances, not yet discovered, which would be observed in diffractive excitation but only weakly couplet to two body formation and decay channels. The SU(3) partners of the Roper resonance N(1470), might be such states and be found with hyperon beams.

The availability of hyperon beams raises the possibility of observing new strong-interaction phenomena not previously available to experiment. The presence of a strange baryon in the initial state allows the study of strange-baryon transitions without strangeness exchange. The present discussion considers three types of strangeness-conserving hyperon transitions that appear to be of interest: (1) electromagnetic transitions, (2) hadron reactions with exchanges of nonstrange Reggeons (ρ, ω, etc.), and (3) diffractive excitation. Electromagnetic transitions can be studied by the

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Primakoff effect, for which a strong excitation of the $\frac{3}{2}^+$ decuplet is expected. However, a U-spin selection rule[1] forbids this excitation for negatively charged hyperons but allows it for neutral and positive baryons. The extent to which this selection rule is violated by SU(3) symmetry breaking is of particular interest, since the most readily available hyperon beams have negative charge. The couplings of nonstrange bosons and Reggeons to strange baryons is of interest because of still untested quark-model and symmetry predictions for these couplings. Diffractive excitation might produce new hyperon resonances which are not excited by strangeness exchange and which appear only weakly in phase-shift analyses—e.g., the SU(3) partners of the Roper N(1470) and other diffractively excited nonstrange baryon resonances.

Theoretical understanding of these questions is not very well founded, but the quark model seems to give a good description of hadron systematics and spectroscopy. These points will therefore be examined with the aid of the quark model to see if any new insight can be obtained.

I. ELECTROMAGNETIC EXCITATION OF BARYON RESONANCES

Decuplet baryons can be excited by the electromagnetic reactions

$$\gamma + p \rightarrow \Delta^+, \tag{1a}$$

$$\gamma + \Sigma^+ \rightarrow Y^{*+}, \tag{1b}$$

$$\gamma + \Sigma^- \rightarrow Y^{*-}, \tag{1c}$$

$$\gamma + \Xi^- \rightarrow \Xi^{*-}, \tag{1d}$$

which can be studied by scattering a hyperon beam on a heavy nucleus and looking for the Primakoff effect[2]

$$\Sigma^{\pm} + Pb \rightarrow Y^{*\pm} + Pb, \tag{2a}$$

$$\Xi^- + Pb \to \Xi^{*-} + Pb. \tag{2b}$$

The excitation of the positively charged baryons—reaction (1a) or (1b)—is allowed by U spin, and the photoproduction process $\gamma + p \to \Delta^+$ is known to be strong. The excitation of the <u>negatively</u> charged hyperons—reaction (1c) or (1d)—is forbidden by U spin. However U-spin conservation is known to be broken. The quark model provides an estimate of the magnitude of this symmetry-breaking effect and also gives a physical picture that guides the intuition better than the abstract algebraic statement that the initial state has $U = \frac{1}{2}$ for all four reactions (1) while the final state has $U = \frac{1}{2}$ for the positive-charge reactions (1a) and (1b) but has $U = \frac{3}{2}$ for the negative-charge reactions (1c) and (1d).

In the quark model, the proton consists of two p quarks coupled to total spin $S_p = 1$ and one n quark which is coupled to the p quarks to give a total spin of $\frac{1}{2}$. In the Δ^+, the three quarks are coupled to a total spin of $\frac{3}{2}$. The electromagnetic excitation of the Δ^+ is found experimentally to be dominated by the magnetic dipole transition.[3] This is predicted by the quark model and has a simple physical interpretation. A magnetic field acting on the quarks in the proton can flip the spin of the n quark relative to that of the p quarks to induce the transitions from the $\frac{1}{2}^+$ to the $\frac{3}{2}^+$ state. Since the charge of the p quark is $Q_p = +\frac{2}{3}$ while the charge of the n quark is $Q_n = -\frac{1}{3}$, the magnetic moments of the p and n quarks have opposite signs. A magnetic field acting on a system of p and n quarks thus rotates the two spins in opposite directions and easily rotates the spin of the n quark relative to that of the p quarks to make the transition. The same picture applies to the excitation of the Σ^+ which contains two p quarks and a λ quark since $Q_\lambda = Q_n = -\frac{1}{3}$.

The negatively charged hyperons Σ^- and Ξ^- contain only n and λ quarks. The Σ^- has two n quarks and one λ; the Ξ^- has two λ quarks and one n. The U-spin selection rule has a simple interpretation in this model. Since the charges of the n and λ quarks are equal, their magnetic moments are also equal in the U-spin symmetry limit. In this limit, an external magnetic field applied to the hyperon rotates all the quarks together by the same amount and therefore leaves the total spin unchanged. There is no force that rotates the spin of the odd quark in a direction opposite to the spin of the other two quarks as in the case of the positive baryons. Thus the dominant transition in the case of negative hyperons is an elastic spin flip of the whole hyperon.

The selection rule is broken if the magnetic moments of the n and λ quarks are not quite equal. In this case the magnetic force on the two types of quarks is not exactly the same and one is rotated more than the other by an external magnetic field. A quantitative estimate of this effect is obtained from the explicit form of the magnetic-moment operator for the baryons in the reactions (1). Each of these baryons contains two identical quarks and one odd quark of a different type. We therefore consider the case of a system containing two quarks of type a and one of type b, where (a, b) = (p, n) for the proton, (p, λ) for the Σ^+, (n,λ) for the Σ^-, and (λ, n) for the Ξ^-.

The magnetic-moment operator $\vec{\mu}$ for a quark is proportional to its spin $\vec{s}$ with proportionality factor g which depends upon the type of quark. Thus

$$\vec{\mu}_a = g_a \vec{s}_a , \tag{3a}$$

$$\vec{\mu}_b = g_b \vec{s}_b . \tag{3b}$$

The total magnetic moment of the three-quark system is then given by

$$\vec{\mu}_B = \sum_i \vec{\mu}_i = g_a \vec{S}_a + g_b \vec{s}_b \tag{4a}$$

where $\vec{S}_a$ is the total spin of the two a quarks. This is conveniently rewritten

$$\vec{\mu}_B = \tfrac{1}{2}(g_a + g_b)\,(\vec{S}_a + \vec{s}_b) + \tfrac{1}{2}(g_a - g_b)\,(\vec{S}_a - \vec{s}_b)$$
$$= \tfrac{1}{2}(g_a + g_b)\,\vec{J}_B + \tfrac{1}{2}(g_a - g_b)\,(\vec{S}_a - \vec{s}_b), \qquad (4b)$$

where

$$\vec{J}_B = \vec{S}_a + \vec{s}_b \qquad (4c)$$

is the total angular momentum of the baryon.

The first term on the right-hand side of Eq. (4b) is proportional to the total-angular-momentum operator and simply rotates the spin of the whole baryon without producing excitations. The entire contribution to the excitations comes from the second term, which is proportional to $g_a - g_b$. This term vanishes, as expected, if the two types of quarks present in the system have equal magnetic moments.

The transition matrix elements for the reactions (1) are proportional to the matrix elements of the second term in Eq. (4b) between the initial and final baryon states. In each of these baryons, two quarks of a type a are coupled to spin 1 and this pair is coupled to one quark of type b to make a spin of $\frac{1}{2}$ in the initial state and a spin of $\frac{3}{2}$ in the final state, thus the calculation of the matrix element is identical in all four cases, except for the multiplicative factor $g_a - g_b$. The cross section is proportional to the square of the transition matrix element. That is,

$$\frac{\bar{\sigma}(\gamma\Sigma^- \to Y^{*-})}{\bar{\sigma}(\gamma p \to \Delta^+)} = \frac{(g_n - g_\lambda)^2}{(g_p - g_n)^2} = \frac{1}{9}\left(1 - \frac{g_\lambda}{g_n}\right)^2, \qquad (5)$$

where $\bar{\sigma}$ denotes the cross section corrected for differences in phase-space factors. In writing this expression, we have substituted

$(a, b) = (n, \lambda)$ for the Σ^- and $(a, b) = (p, n)$ for the proton and we have used the quark-model value $g_p = -2g_n$ that gives excellent agreement for the ratio of neutron and proton magnetic moments.

The result (5) indicates that U-spin symmetry breaking should have a very small effect on the breaking of the selection rule. Even if g_λ differs from g_n by 30% (which would be an appreciable symmetry breaking), the Σ^- excitation cross section would still be only 1% of that for the Δ.

Another possible breaking of SU(3) symmetry is an admixture of decuplet in the Σ^-, analogous to the d-wave admixture in the deuteron. Estimates of such admixtures give 15% or 20% in amplitude.[4] Their contribution to transition probabilities is proportional to the square of the mixing amplitude. This again suggests effects of only a few percent. An exact calculation of this effect is more complicated and not feasible. There are too many unknown parameters, since this transition would be electric quadrupole rather than magnetic dipole and there could be mixing in the Y^* as well as in the Σ^-.

II. NEW COUPLINGS OF HYPERONS AND BOSONS

All this theoretical argument can be questioned because of experimental evidence that quark-model predictions fail for the production of Σ and Y^* hyperons in strangeness-exchange reactions on nucleons.[5] This suggests that the naive calculations of the three-point couplings of vector mesons between nonstrange octet baryons and strange decuplet hyperons is not given properly by the quark model. However the couplings of the nonstrange vector mesons to nonstrange baryons are well described by the model, as indicated by the success of the quark model in predicting[2] nucleon magnetic moments and total cross sections for reactions on nucleon targets and for the Δ photoproduction reaction (1a). The hyperon transitions (1c) and (1d) test the coupling of a nonstrange vector boson between strange octet and strange decuplet baryons. So far this particular coupling has not been tested. It is therefore of considerable theoretical interest to test the U-spin selection rule. This can be done by looking for Primakoff excitation of negative hyperons from heavy nuclei, as in the reactions (2), and comparing the cross sections with the corresponding Primakoff excitation of the Δ from a nucleon beam to test the relation (5).

The quite interesting couplings of nonstrange bosons or Regge trajectories to strange octet baryons are also not yet tested. These can be investigated by examining the hyperon-baryon total cross sections. There are definite quark-model and symmetry predictions that relate these hyperon cross sections to meson-nucleon and nucleon-nucleon total cross sections that have already been observed. The quark model predicts that the Σ^-N and Ξ^-N total cross sections should satisfy the relations[6, 7]

$$\sigma(pn) - \sigma(\Sigma^-p) = \sigma(K^-n) - \sigma(\pi^+p) \approx 4 \text{ mb}, \tag{6a}$$

$$\sigma(\Sigma^-p) - \sigma(\Xi^-p) = \sigma(K^-n) - \sigma(\pi^+p) \approx 4 \text{ mb}, \tag{6b}$$

where the value of 4 mb, which is nearly independent of energy, is taken from experiment.[7] Thus $\sigma(\Sigma^-p)$ and $\sigma(\Xi^-p)$ are expected to be nearly independent of energy and below the nucleon-nucleon cross sections by about 4 mb and 8 mb, respectively.

Regge descriptions give similar predictions. The Σ^-N and Ξ^-N cross sections are predicted to be similar in character to NN total cross sections--i.e., they are dominated by the Pomeron. The contributions of the leading exchange-degenerate secondary Regge trajectories are expected to cancel in these exotic ($B = 2$) channels.[8] In the SU(3) symmetry limit, the Pomeron contributions to all baryon-baryon cross sections are expected to be equal. In meson-baryon cross sections, this symmetry is seen to be broken; the KN Pomeron contribution seems to be about 4 mb below the πN contribution. The quark-model predictions (6) can be interpreted as requiring a similar symmetry breaking for the Pomeron contributions in meson-baryon and baryon-baryon scattering.

If the predictions (6) are verified by experiment, this will constitute further evidence for exchange degeneracy of the leading trajectories and SU(3) symmetry breaking in the Pomeron couplings. Values for the couplings of the leading Regge trajectories are then not obtained from these cross sections because of the cancellations of exchange degeneracy. They will be obtained from antihyperon-nucleon cross sections, in which the contributions from the exchange-degenerate pairs add rather than cancel.

When antihyperon total cross sections become available, it will be interesting to examine linear combinations of cross sections that project out t-channel exchanges with definite quantum numbers. The odd-signature isoscalar exchange (ω exchange) is of particular interest, since predictions from omega universality[9] have so far been in remarkable agreement with experimental data for kaon-nucleon and nucleon-nucleon total cross sections, but have not yet been tested for strange baryons. One omega-universality prediction is

$$\sigma(\bar{\Lambda}p) - \sigma(\Lambda p) = \frac{2}{3}\,[\sigma(\bar{p}p) - \sigma(pp)] - \frac{1}{3}\,[\sigma(\pi^- p) - \sigma(\pi^+ p)]. \qquad (6c)$$

For the even-signature isoscalar exchange (Pomeron and f^0 exchange), the quark model predicts

$$\sigma(\bar{\Lambda}p) + \sigma(\Lambda p) = \tfrac{1}{2}[\sigma(\pi^- p) + \sigma(\pi^+ p) + \sigma(K^- p) + \sigma(K^+ p) + \sigma(K^- n) + \sigma(K^+ n)]. \quad (6d)$$

The corresponding quark-model prediction for even-signature isoscalar exchange for nonstrange baryons disagrees with experiment. The ratio of baryon-baryon to meson-baryon cross sections is 15—20% greater than the famous quark-model prediction of $\frac{3}{2}$ and seems to be independent of energy.[10] It will be interesting to see whether a similar discrepancy prevails for the relation (6d).

III. DIFFRACTIVE EXCITATION OF NEW RESONANCES

Diffractive excitation of hyperons might reveal new resonances not previously discovered; for example, there might be strange partners to the Roper (1470) nucleon resonance. A very simple picture can be used in the quark model to explain how certain resonances would be excited diffractively and would not be seen in common excitation mechanisms. In a variety of approaches, the coupling of baryon resonances to meson-baryon channels is described as due to a transition of a single quark in the baryon while the other two quarks remain as spectators.[11] This is true (1) in the model of boson emission as the emission of a boson quantum by a single quark,[12] (2) in the Harari-Rosner duality-diagram description of production and decay of resonances,[13] and (3) in the Levin-Frankfurt model[6] of meson-baryon scattering.

In a three-body system, there are two independent orbital angular momenta. In resonance excitation with this simple "spectator" picture, it is convenient to choose these two orbital angular momenta to be (1) the relative orbital angular momentum of the two spectator quarks and (2) the angular momentum of the active quark with respect to the center of mass of the spectator quarks. The states of the spectator quarks and their relative orbital angular momentum do not change in such a transition. In the symmetric quark model, the quark in the low-lying baryon octet and decuplet are all in relative s states. Thus this model suggests that the only resonances excited in this way are those in which two quarks remain in a relative s state and the total orbital momentum is carried by the third quark. States that involve simultaneous excitation of both orbital angular momenta are not reached by this mechanism.[11]

Difficulties arise in the application of this picture to diffraction dissociation processes. These are most simply seen in the case of A1 and Q production by pion and kaon beams, respectively.

In the quark model, these excitations are $0^- \to 1^+$ transitions between the $^1S_0$ and $^3P_1$ configurations and require both a spin flip and an orbital excitation by one unit of orbital angular momentum. If this is produced by single-quark scattering, the orbital and spin contributions would not be expected to be correlated. Thus, a final state with total spin S = 1 and orbital angular momentum L = 1 would not have L and S coupled to a definite total angular momentum J, but would be a mixture of 0^+, 1^+, and 2^+ states. This picture predicts excitation of all three states, with ratios depending upon Clebsch-Gordan coefficients, in disagreement with Morrison's rule[14] that only unnatural-parity 1^+ states are excited by diffraction dissociation—the natural-parity 0^+ and 2^+ states are not.

The undesired excitation of 0^+ and 2^+ states can be eliminated by adjusting the relative phases of the two terms arising in the model from quark and antiquark excitation so that they cancel one another in the 0^+ and 2^+ excitations and add constructively in the 1^+ excitation. This can be done elegantly by writing the expression for the transition matrix element of the hadron undergoing diffraction dissociation in the form

$$\langle B \,|\, T(\vec{q}) \,|\, A\rangle = \langle B \,|\, \sum (\vec{\sigma}_i \times \vec{p}_i \cdot \vec{q})\, e^{i\vec{q}\cdot\vec{r}_i} \,|A\rangle\,, \tag{7}$$

where A and B are the initial and final hadron states, σ_i, $\vec{p}_i$, and $\vec{r}_i$ are the spin, momentum, and coordinate of the ith quark, $\vec{q}$ is the momentum transfer, and the sum is over all quarks and antiquarks. The operator appearing in the matrix element is a sum of single-quark operators, as required by the simple picture. The leading term in the expansion of the exponential is a polar vector and can only connect a 0^- initial state to a 1^+ final state. If this result (7) can be derived from some dynamical principle, it would give an elegant derivation of Morrison's rule. However, such a single-quark operator is rather artificial in the simple picture.

We now consider a new alternative description for diffraction dissociation, in the spirit of "two-component duality" in which diffractive processes that remain constant at high energy are qualitatively different from other exchanges that decrease with energy.[15] It also fits the general philosophy that diffraction dissociation is a kind of elastic process in which the probability of exciting a particular resonance depends on the overlap of the initial- and final-state wave functions. In this description of diffraction, elastic scattering dominates at low momentum transfer because all excited-state wave functions are orthogonal to the initial state and have no overlap. However, at high momentum transfers the initial and final baryons are not at rest in the same Lorentz frame. Therefore the overlap integral calculated, say, in the rest frame of the produced resonance includes a Lorentz boost of the initial baryon state. This changes the wave function by producing Lorentz contractions and Wigner rotations of the spins and thus introduces components that can look like isobars and produce inelastic transitions. Because of the mass differences, there is always a considerable momentum transfer in isobar production even in the forward direction, and a Lorentz boost is always required.

Since the Lorentz boost affects all of the three quarks in the baryon, there is no reason to prefer excitations of only a single quark from the initial state. Diffractive excitation can thus produce states that differ from the target by simultaneous excitation of two quarks. Such states would decay dominantly via modes that are not quasi-two-body, since the decay to the baryon octet would require de-excitation of two orbital angular momenta and hence would be forbidden by the common mechanism. A $\frac{1}{2}^+$ state could be formed by having each of the two internal orbital angular momenta excited to a p wave since these two 1-unit orbital angular momenta can be coupled with the three quark spins to give a total angular momentum of $\frac{1}{2}$. One can see qualitatively how such an excitation occurs in a simple model analogous to atomic physics, with quarks described by solutions of the Dirac equation in an external field. The lowest state of the hydrogen atom is an $s_{1/2}$ state but there are small components in the wave function that look like $p_{1/2}$. A Lorentz boost mixes the $s_{1/2}$ and the $p_{1/2}$ wave functions. Such a transformation on two quarks adds two units of orbital angular momentum while leaving the total angular momentum j of each individual quark unchanged. This oversimplified picture should not be taken too literally, but the observation that baryon resonances that are not coupled strongly to quasi-two-body channels might be diffractively produced may well have a more general validity.

This model for diffractive processes can be formulated more quantitatively as follows. Hadron wave functions from the nonrelativistic quark model are considered to be reasonable approximations only in the hadron rest frame.[16] For hadrons in motion, there are Lorentz contractions of the spatial wave functions and Wigner rotations of the spins. We assume that as a result of these relativistic effects, the wave functions of states produced by diffraction dissociation are mixed into the wave function of the incident hadron. We assume that in the diffraction process the hadron receives a momentum transfer without any change in the other degrees of freedom. This is consistent with the philosphy that the Pomeron carries the quantum numbers of the vacuum. Thus if a quark in a hadron is described by a Dirac spinor, the transition is described by multiplying the wave function of each quark by a plane wave that does not affect the spinor degrees of freedom. The result is

$$\langle B \mid T_D(\vec{q}) \mid A \rangle = \langle B \mid \prod_i e^{i\vec{q}_i \cdot \vec{r}_i} \mid A \rangle , \tag{8}$$

where the subscript D (diffraction) distinguishes the expression (8) from (7), and $\vec{q}_i$ is the portion of the momentum transfer given to the ith quark.

For convenience consider the process in the rest frame of the final state. This is also the Jackson frame commonly used for the analysis of the decay angular distributions of the resonances produced. The wave function of the final state at rest is given directly by the nonrelativistic quark model. The initial-state wave function that describes a moving hadron is constructed by boosting the nonrelativistic quark-model wave function from the rest frame to the required incident velocity. That is,

$$|A(\vec{v})\rangle = U_{A\vec{v}}\,|A\,(0)\rangle\,, \tag{9}$$

where $|A(0)\rangle$ and $|A(\vec{v})\rangle$ are the wave functions of the initial state in its rest frame and in the Jackson frame, respectively, and $U_{A\vec{v}}$ is the Lorentz boost from the rest frame of the initial hadron to the state in which it is moving with velocity $\vec{v}$. The boost $U_{A\vec{v}}$ acts both on the spatial degrees of freedom and on the Dirac spinor indices.

If the motion of the quarks within the hadron is considered as nonrelativistic, the spatial part of the boost simply shifts the momentum by multiplying the wave function by a plane wave. We neglect the Lorentz contraction of the wave function. Since this boost brings the wave function from rest to the incident velocity while the transition from the initial to final states (2) brings the hadron to rest, the spatial part of the Lorentz boost in Eq. (3) cancels the plane wave in the transition matrix element (2) if the momentum transfer is appropriately divided between the constituent quarks. If these spatial factors do not exactly cancel the spatial form factor that remains will be neglected at this stage together with such other spatial effects as the Lorentz contraction.

Equation (9) is now used to write a simple expression for the transition (8) in which all wave functions are expressed in their rest frames and can be described by the nonrelativistic quark model. Since the spatial effects cancel in our approximation, only spin effects remain and

$$\langle B\,(0)\,|\,T_D\,|\,A(\vec{v})\rangle = \langle B(0)\,|U^{(s)}_{A\vec{v}}\,|\,A(0)\rangle\,, \tag{10}$$

where A(0) and B(0) are the quark-model hadron wave functions in the respective rest frames and $U^{(s)}_{A\vec{v}}$ is the spin part of the Lorentz boost that takes the incident particle A from its rest frame to the velocity $\vec{v}$ in the Jackson frame.

Two qualitative features are immediately evident from this expression (10).

1. Morrison's rule. The operator $U^{(s)}_{A\overline{V}}$ is an exponential function of the generator of Lorentz transformations. Since the generator is a polar vector, the expansion of the exponential in a power series contains different powers of the same polar vector—i.e., only terms that have natural parity. If the initial state A is a pseudoscalar meson, there will be nonvanishing excitation only to states having unnatural parity 0^-, 1^+, 2^-, . . . as given by Morrison's rule. For baryons the situation is more complicated because the initial spin of $\frac{1}{2}$ can always be coupled in two ways to a given natural-parity excitation.

2. Deviations from the simple spectator picture. The transition operator is manifestly not a single-quark operator since it produces Lorentz transformations on all the quarks in the system. This model thus predicts that diffraction dissociation can excite states that are not produced by conventional exchanges. This might explain the apparent anomaly associated with the N(1470).

These points are somewhat illuminated by examining the expansion of the Lorentz transformation in powers of its generator M. In this expansion

$$e^{ivM} = 1 + ivM + \tfrac{1}{2}(iv)^2M^2 + \cdot \cdot \cdot , \tag{11}$$

the Lorentz generator M is an "odd" operator in the Dirac sense. In the nonrelativistic limit, it connects "large" and "small" components of a Dirac wave function. Thus in a nonrelativistic approximation, the matrix element of the operator (11) for a given transition should be given by the leading nonvanishing term in the expansion. The first term describes elastic scattering. Since M is a polar vector, the second term describes the excitation of states whose parity is opposite that of the ground state and the third term describes excitations with the same parity as the ground state.

The second term (which describes the odd-parity excitations) can still satisfy the simple spectator picture under the reasonable assumption that the Lorentz generator M is additive in the individual quark variables to a good approximation. However, the third term (which describes the even-parity excitation) is manifestly not additive and contains two-quark excitations even if M is additive. With this interpretation, production of even-parity states with two-quark excitations should be expected in diffraction dissociation. Such two-quark excitations would not be produced by nondiffractive exchanges described by the simple picture and would not decay by the conventional quark-model mechanism to a meson-baryon two-body channel. They could not be observed as resonances in meson-baryon phase-shift analyses. This would support Morrison's conjecture that resonances produced diffractively are somehow different from resonances produced in other ways.[14] Furthermore, since there are as yet no diffractively produced hyperon resonances, the SU(3) partners of diffractively produced double-quark excitations would not yet have been observed, and there would be difficulties in classifying such resonances as the N(1470) in SU(3) representations.

Quantitative calculations of diffraction dissociation in this model are more difficult because they depend more critically on effects that are normally neglected in the nonrelativistic quark model. Since the generator M is an odd operator, it necessitates the use of the small components of the wave function as well as the large components. Conventional nonrelativistic reduction procedures used in nuclear physics are ambiguous in the quark model. The use of the Dirac equation to connect small and large components requires some knowledge of the interaction potential and, in particular, of whether it is a four-vector or world scalar. These ambiguities become particularly significant and annoying for the case of very strong binding.

Even the most naive calculations can only be expressed in terms of a free parameter having the dimensions of a mass and a value anywhere between 1/3 the baryon mass and the mass of a free quark.

The picture of a Dirac particle moving in some kind of self-consistent central field (the analog of the one used for relativistic corrections in nuclear physics) does provide some intuition for qualitative thinking about this model. In the symmetric quark model, the three quarks in a baryon are considered to be in $s_{1/2}$ orbits. In the Dirac description, these would have small components that look like $p_{1/2}$ wave functions. A Lorentz transformation acting on this spinor mixes the $s_{1/2}$ and $p_{1/2}$ wave functions and thus can produce configurations in which the quark is excited up to the $p_{1/2}$ orbit. A single-quark excitation of this type describes the excitation of the 1^+ meson resonances A1 and Q and of $\frac{1}{2}^-$ and $\frac{3}{2}^-$ baryon resonances that are p-wave excitations in the quark model. Such excitations should also appear in the Levin-Frankfurt model since they are single-quark excitations. The next even-parity excitations would then be described by simultaneously exciting two quarks to $p_{1/2}$ orbits. This gives states that are not excited in the Levin-Frankfurt model. The first example of such a state might be the N(1470).

Excitations of higher resonances require higher terms in the expansion (5) and cannot be considered in a consistent way without re-examining the entire model from the beginning and looking for other higher order relativistic effects. These include, for example, the Lorentz contraction of the spatial part of the wave function produced by the boost $U_{a\vec{v}}$ in Eq. (9). This has been neglected in the derivation of Eq. (10) by making the nonrelativistic approximation, in which all momentum components in the wave function are shifted by the same amount $\vec{q}$ in the Lorentz transformation. Taking this effect into account introduces an additional orbital factor in Eq. (10), and it again would produce only natural-parity excitations but would allow for orbital excitations of quarks beyond the $s_{1/2}$ and $p_{1/2}$ orbits. However, the precise nature of excitations seems to depend on detailed characteristics of the wave functions and cannot be described as reliably as the qualitative features.

The principal experimentally verifiable qualitative predictions of this model are baryon resonances that are observed only in diffraction dissociation and whose decay into two-body channels is strongly inhibited. With hyperon beams available for diffraction dissociation studies of strange baryons, such resonances should be observed as the previously missing SU(3) partners of diffractively excited nonstrange baryon resonances. The first such resonances to be expected would be the SU(3) partners of the N^* (1470), which would be a Y^*(1620) and a Ξ^*(1770).

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